

Higher Derivative D-term Inflation in New-minimal Supergravity

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Abstract

We revisit the D-term inflation and amend it with ghost-free higher derivative couplings of chiral superfields to super-curvature. These couplings realize a more generic inflationary phase in supergravity. After pointing out that a consistent embedding of these specific higher derivatives is known to exist only in the new-minimal supergravity, we show how a potential for the scalar component may arise due to a Fayet-Iliopoulos D-term. We then turn to inflationary cosmology and explicitly discuss different types of potentials, which capture properties of the common scenarios. These models thanks to the derivative coupling: i) naturally evade the supergravity η -problem, ii) drive inflation for a wider range of parameter values, and iii) predict lower values for the tensor-to-scalar ratio.

1 Introduction

Observational data strongly indicate that an inflationary phase did occur at some stage in the early universe. Either in supersymmetric or non-supersymmetric theories the slow-roll inflation is the dominant paradigm [1, 2, 3]. This phase is characterized by the Hubble friction, hence theories that generate enhanced friction effects are cosmologically rather motivated. It has been found that when a scalar field has derivative couplings to curvature, then it can slow-roll down even at relatively steep potentials during a (nearly) de Sitter phase. Nevertheless, not all derivative couplings to curvature are consistent, but there exist specific classes which lead to viable field theories [4]. An example of these ghost-free higher derivatives is the kinetic coupling of a scalar field to the Einstein tensor

$$\frac{1}{M_*^2} G^{mn} \partial_m \phi \partial_n \phi \quad (1)$$

which has given rise to a considerable amount of scientific activity [5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16, 17]. In fact there also exist even higher order consistent derivative couplings [4, 18, 19, 20, 21].

The coupling (1) has been called “Gravitationally Enhanced Friction” (GEF) mechanism [12]. The attraction of the mechanism is that it can set more general initial conditions for the inflationary phase, by relaxing the slow-roll conditions. Note that the mass scale in (1) has to be smaller than the Hubble scale during inflation because only if $M_* \ll H$ the enhanced friction effects are more influential and noticeable. This raises the question of the origin of this scale; it is rather motivating to find it among the dilaton couplings of the heterotic string [22].

Thus, it would be desirable to have an *embedding* of the GEF mechanism in the theory of supergravity, which is the framework under which supersymmetric theories during inflation should be studied. In a supergravity theory, which is a non-renormalizable theory, the number of couplings that need to be specified is in principle infinite. The cut-off of the theory is considered to be the Planck mass, M_P , and an estimation in supergravity ceases to be reliable for field values $\phi \gg M_P$, unless the non-renormalizable terms are suppressed, e.g. due to a particular symmetry. Even though one can construct models in which the inflaton field value experiences sub-Planckian variation, a generic supergravity theory will fail to inflate because the F-term part of the potential yields a too big inflaton mass [3].

The embedding of the GEF mechanism is in general not a trivial task; as we have mentioned, higher derivative theories may come accompanied by ghost instabilities. In earlier works [23, 24] it was understood, that in order for this coupling to be consistently realized, one has to turn to the new-minimal supergravity [25, 26, 27]. Nevertheless, the coupling is still inconsistent unless the non-minimally coupled superfield has a vanishing R-charge, and is neutral under any gauge group. Hence it is not possible to endow this superfield with a conventional self-interaction in terms of superpotential or gauging. This problem can be solved, as we propose here, by breaking supersymmetry with a Fayet-Iliopoulos term, which may induce a potential for the scalar field.

In this work we explore the modified dynamics of a non-minimally coupled superfield to curvature, and we find that inflation can be realized and described reliably in a supergravitational framework. Indeed, the scalar component of the Φ superfield is governed solely by a D-term potential and experiences a high friction during a de Sitter phase. Moreover, due the vanishing of the superpotential, it is also expected to have negligible interactions with other fields, a fact that is supported by the Planck observational data. Therefore the superfield Φ is tailor-made for driving inflation in supergravity.

The motivation of this article is both particle theoretical and cosmological. On technical grounds, in section 2, we show how it is possible to introduce a potential for the non-minimally coupled field when it is coupled to supergravity. Then, in section 3, we show how an inflating theory driven by the supersymmetric slotheon - it has been named slotheon after [28], can evade some shortcomings common in conventional inflationary supergravity and we revisit particular inflationary models. We find the new field space region where an accelerated expansion is realized and check whether these models can fit the observational data even though they were previously excluded.

2 New-minimal supergravity: Derivative couplings and D-terms

The minimal theories of supergravity have a rich structure originating from the possible compensating multiplets that break the underlying superconformal theory to super-Poincaré [29, 30]. The underlying dualities among the compensating multiplets survive the gauge fixing and lead to equivalent couplings to matter [31], but break down as soon as higher derivatives are introduced. The couplings we want to study here make this duality-breakdown even more manifest, since the only known supergravity which

can accommodate them in a consistent way [23, 24] is the so-called *new-minimal supergravity* [25, 26, 27]. An aspect of the new-minimal supergravity not encountered in the *standard* supergravity is the necessary existence of a chiral symmetry. It is well known that rigid supersymmetry allows for the existence of a *chiral symmetry* called $U(1)_R$. This *R-symmetry* becomes local and is gauged by one of the auxiliary fields of the gravitational supermultiplet.

The new-minimal supergravity [25] is the supersymmetric theory of the gravitational multiplet

$$e_m^a, \quad \psi_m^\alpha, \quad A_m, \quad B_{mn}. \quad (2)$$

The first two fields are the vierbein and its superpartner the gravitino, a spin- $\frac{3}{2}$ Rarita-Schwinger field. The last two fields are auxiliaries. The real auxiliary vector A_m gauges the $U(1)_R$ chiral symmetry. The auxiliary B_{mn} is a real two-form appearing only through its dual field strength H_m , which satisfies $\hat{D}^a H_a = 0$, for the supercovariant derivative \hat{D}^a . This constraint can be solved in terms of B_{mn} . Note that all the fields of the new-minimal supergravity multiplet are gauge fields.

We will use superspace techniques to guarantee that our component form Lagrangians are supersymmetric. The interested reader may consult for example [27] where a treatment of the new-minimal superspace is given. The new minimal supergravity free Lagrangian is given by

$$\mathcal{L}_{\text{sugra}} = -2M_P^2 \int d^4\theta E V_R. \quad (3)$$

Here V_R is the gauge multiplet of the R-symmetry, which (in the appropriate WZ gauge) contains the auxiliary fields in its vector component, $-\frac{1}{2}[\nabla_\alpha, \bar{\nabla}_{\dot{\alpha}}]V_R| = A_{\alpha\dot{\alpha}} - 3H_{\alpha\dot{\alpha}}$, and the Ricci scalar in its highest components, $\frac{1}{8}\nabla^\alpha \bar{\nabla}^2 \nabla_\alpha V_R| = -\frac{1}{2}(R + 6H^a H_a)$. The E is the super-determinant of new-minimal supergravity, but in general (as we also do here) one can calculate the supersymmetric Lagrangians only with the use of the F-term formula, since

$$\int d^4\theta EX = \frac{1}{2} \int d^2\theta \mathcal{E} \left(-\frac{1}{4} \bar{\nabla}^2 X \right) + c.c. \quad (4)$$

In the chiral *theta* expansion the chiral density reads $\mathcal{E} = e + ie\sqrt{2}\theta\sigma^a\bar{\psi}_a - \theta^2 e\bar{\psi}_a\bar{\sigma}^{ab}\bar{\psi}_b$. Note that X is a generic hermitian superfield with vanishing chiral weight, while its chiral projection $(-\frac{1}{4}\bar{\nabla}^2 X)$ has chiral weight $n = 1$. The bosonic sector of Lagrangian (3) is

$$\mathcal{L}_{\text{sugra}}^B = M_P^2 e \left(\frac{1}{2}R + 2A_a H^a - 3H_a H^a \right). \quad (5)$$

For the matter sector we have a chiral multiplet, defined by $\bar{\nabla}_{\dot{\alpha}}\Phi = 0$ which has bosonic components, a physical complex scalar $A = \phi + i\beta$, and an auxiliary complex field F , defined as

$$\Phi| = A, \quad -\frac{1}{4}\nabla^2\Phi| = F. \quad (6)$$

In general, a chiral superfield in new-minimal supergravity is allowed to have an arbitrary R-charge, but we stress that our chiral superfield has a vanishing one [23] in order to avoid ghost instabilities

$$n_\Phi = 0. \quad (7)$$

The minimal kinematic Lagrangian for this multiplet is in superspace

$$\mathcal{L}_0 = \int d^4\theta E \bar{\Phi}\Phi \quad (8)$$

the bosonic sector of which is

$$\mathcal{L}_0^B = A\Box\bar{A} + F\bar{F} - iH^m (A\partial_m\bar{A} - \bar{A}\partial_m A). \quad (9)$$

Finally, concerning our chiral superfield, it will also have a non-minimal derivative coupling with the supergravity multiplet

$$\mathcal{L}_{M_*} = iM_*^{-2} \int d^4\theta E [\bar{\Phi} E^a \nabla_a \Phi] + c.c. \quad (10)$$

where E^a is a curvature real linear superfield ($\nabla^2 E_a = \bar{\nabla}^2 E_a = 0$) of the new-minimal supergravity. The E^a superfield has bosonic components $E_a| = H_a$ and $\frac{1}{4}\bar{\sigma}_a^{\dot{\alpha}\alpha}[\nabla_\alpha, \bar{\nabla}_{\dot{\alpha}}]E_b| = \frac{1}{2}(G_{ab} - g_{ab}H^c H_c - 2H_a H_b - {}^*\mathcal{F}_{ab})$, where $G_{mn} = R_{mn} - \frac{1}{2}g_{mn}R$ is the Einstein tensor and $\mathcal{F}_{mn} = \partial_m A_n - \partial_n A_m$ is the field strength of the supergravity auxiliary field A_m . It is worth mentioning that E^a satisfies the superspace Bianchi identity $\nabla^a E_a = 0$. For a discussion and derivation of the Lagrangian (10) see [23]. The bosonic sector of Lagrangian (10) is

$$\begin{aligned} \mathcal{L}_{M_*}^B = & M_*^{-2} \left[G^{ab} \partial_b \bar{A} \partial_a A + 2F\bar{F}H^a A_a - 2F\bar{F}H^a H_a \right. \\ & + iH^a (\bar{F}\partial_a F - F\partial_a \bar{F}) - \partial_b A \partial^b \bar{A} H_a H^a \\ & \left. + 2H^a \partial_a A H^b \partial_b \bar{A} - iH_c (\partial_b \bar{A} \mathcal{D}^c \partial^b A - \partial_b A \mathcal{D}^c \partial^b \bar{A}) \right]. \end{aligned} \quad (11)$$

Note that this term, although it contains higher derivatives, does not lead to ghost states or instabilities. The ghost instabilities are in fact evaded due to the vanishing chiral weight of the chiral superfield Φ . On the other hand, the vanishing chiral weight forbids the self-coupling via a superpotential due to the R-symmetry. Thus this superfield is not allowed to have a superpotential. Moreover it is also not allowed to be gauged, since this will also give rise to ghost instabilities via inconsistent derivative couplings of the gauge fields to curvature. The only remaining option is the indirect introduction of self-interaction via a *gauge kinetic function*.

The gauge sector of our theory is composed of a standard $U(1)$ gauge multiplet V , with a Φ -dependent gauge kinetic function and a Fayet-Iliopoulos term. The $U(1)$ gauge multiplet consists of a gauge vector field v_m , a majorana gaugino λ^α , and a real auxiliary field D . In particular, the definition of the bosonic components of the vector is

$$-\frac{1}{2}[\nabla_\alpha, \bar{\nabla}_{\dot{\alpha}}]V| = v_{\alpha\dot{\alpha}}, \quad \frac{1}{8}\nabla^\alpha \bar{\nabla}^2 \nabla_\alpha V| = D. \quad (12)$$

Note that due to the structure of new-minimal supergravity, a FI term is in general allowed (even the superspace Lagrangian of pure new-minimal supergravity (3), is a FI term). Thus we have in superspace

$$\mathcal{L}_g = \frac{1}{4} \int d^2\theta \mathcal{E} f(\Phi) W^2(V) + c.c. + 2\xi \int d^4\theta EV \quad (13)$$

with $W_\alpha(V) = -\frac{1}{4}\bar{\nabla}^2 \nabla_\alpha V$. The $f(\Phi)$ is a holomorphic function of the chiral superfield Φ and ξ is the Fayet-Iliopoulos parameter of mass dimension two. The bosonic sector of (13) reads

$$e^{-1}\mathcal{L}_g^B = -\frac{1}{4}\text{Re}f(A)F^{mn}F_{mn} + \frac{1}{4}\text{Im}f(A)F^{mn}{}^*F_{mn} + \frac{1}{2}\text{Re}f(A)D^2 + \xi D - 2\xi v_a H^a \quad (14)$$

where $F_{mn} = \partial_m v_n - \partial_n v_m$. We stress that the scalar A is not charged under this $U(1)$, thus there is no restriction in the form of $f(A)$ apart from holomorphicity.

The total Lagrangian we are interested in is

$$\mathcal{L}_{\text{total}} = M_P^2 \mathcal{L}_{\text{sugra}} + \mathcal{L}_g + \mathcal{L}_0 + \mathcal{L}_{M_*} \quad (15)$$

which reads

$$\begin{aligned} e^{-1} \mathcal{L}_{\text{total}} = & M_P^2 \left[\frac{1}{2} \mathcal{R} + 2V^a H_a - 3H^a H_a \right] + A \square \bar{A} + F \bar{F} \\ & + M_*^{-2} \left[G^{ab} \partial_b \bar{A} \partial_a A - 2F \bar{F} H^a H_a - \partial_b A \partial^b \bar{A} H_a H^a + 2H^a \partial_a A H^b \partial_b \bar{A} \right] \\ & - \frac{1}{4} \text{Re} f(A) F^{mn} F_{mn} + \frac{1}{4} \text{Im} f(A) F^{mn} * F_{mn} \\ & + \frac{1}{2} \text{Re} f(A) D^2 + \xi D, \end{aligned} \quad (16)$$

where we have redefined the auxiliary field A^a to V^a

$$\begin{aligned} V^a = & A^a \left(1 + \frac{1}{M_P^2} M_*^{-2} F \bar{F} \right) - \frac{1}{M_P^2} \xi v^a \\ & + \frac{1}{2M_P^2} \left(i \bar{A} \partial^a A - i A \partial^a \bar{A} - i M_*^{-2} F \partial^a \bar{F} + i M_*^{-2} \bar{F} \partial^a F \right) \\ & + \frac{1}{2M_P^2} \left(i M_*^{-2} \partial_b A \mathcal{D}^a \partial^b \bar{A} - i M_*^{-2} \partial_b \bar{A} \mathcal{D}^a \partial^b A \right). \end{aligned} \quad (17)$$

Lagrangian (16) contains four auxiliary fields. First, by solving the equations of motion for the supergravity auxiliary fields we find that the vector H_m vanishes and V_m reduces to a pure gauge. In fact since H_m is the dual field-strength of B_{mn} here both auxiliary fields of the new-minimal supergravity are pure gauge on-shell. For the auxiliary field F it is easy to see that it will also vanish on-shell. Finally, by solving the equations of motion for the auxiliary field D of the gauge multiplet we find

$$D = -\frac{\xi}{\text{Re} f(A)}. \quad (18)$$

After plugging back our results, we have the following on-shell form for (16)

$$\begin{aligned} e^{-1} \mathcal{L}_{\text{total}} = & \frac{M_P^2}{2} \mathcal{R} + A \square \bar{A} + M_*^{-2} G^{ab} \partial_a \bar{A} \partial_b A - \frac{1}{2} \frac{\xi^2}{\text{Re} f(A)} \\ & - \frac{1}{4} \text{Re} f(A) F^{mn} F_{mn} + \frac{1}{4} \text{Im} f(A) F^{mn} * F_{mn}. \end{aligned} \quad (19)$$

Note that this Lagrangian (19) does *not* contain ghost states or instabilities.

The Fayet-Iliopoulos term inside (14) breaks supersymmetry, and combining it with the gauge kinetic function has the effect of introducing a scalar potential which reads

$$\mathcal{V} = \frac{1}{2} \frac{\xi^2}{\text{Re} f(A)}, \quad (20)$$

where $A = \phi + i\beta$. This is the D-term potential. It is expected that only in the case of broken supersymmetry one can have a potential for the A field which has the non-minimal kinetic coupling to gravity. This stems from the fact that the chiral $U(1)_R$ symmetry of new-minimal supergravity forbids a potential for this field, due to its vanishing chiral weight. The advantage of a Fayet-Iliopoulos term is that it breaks supersymmetry spontaneously.

3 Application to inflation

3.1 A pure D-term inflation

In the standard supergravity the scalar potential of chiral superfields transforming in some representation of a gauge group has the following form¹

$$V = e^{K/M_P^2} \left[F_i (K^{-1})^i_j F^j - 3 \frac{|W|^2}{M_P^2} \right] + \frac{g^2}{2} \frac{1}{\text{Ref}_{ab}} D^a D^b \quad (21)$$

where $F^i = W^i + K^i W/M_P^2$ and $D^a = K^i (T^a)_i^j z_j + \xi^a$. The upper (lower) index i denotes derivatives with respect to the ϕ_i (ϕ^{*i}) field. The slow-roll conditions imply

$$\epsilon \ll 1 \Rightarrow \frac{K_\phi}{M_P} + \dots \ll 1 \quad (22)$$

$$\eta \ll 1 \Rightarrow 3K_{\phi\bar{\phi}} H^2 + \dots \ll H^2. \quad (23)$$

Here the subscript ϕ denotes a derivative with respect to the inflaton. The inflaton vacuum energy dominates the energy density of the universe and the relation $H^2 = V/(3M_P^2)$ has been used in the second condition $\eta \ll 1$. In the low energy minimum the Kähler metric should be normalized to one and it is not expected to be suppressed during inflation. Therefore F-type inflation in supergravity theories is hard to be realized unless the Kähler and the superpotential have a special form or accidental cancellations take place [3, 1, 37].

A resolution to this η -problem in generic supergravity theories can be given by a symmetry that suppresses the F-term part of the scalar potential. In the presence of such a symmetry the potential is naturally dominated by a Fayet-Iliopoulos D-term which exists for $U(1)$ gauge groups. Here, the R-symmetry of the theory forbids the superpotential interactions for the A field non-minimally coupled to the G^{mn} tensor. The spontaneous breaking of supersymmetry during inflation may introduce interactions however these will be generated radiatively and should not affect the tree level D-term inflationary potential. The D-term potential domination together with the enhanced friction features strongly motivates the study of this higher derivative theory to inflationary applications.

3.2 Expanding the allowed initial conditions for inflation

The complex scalar field A is governed by the scalar potential generated by the Fayet-Iliopoulos supersymmetry breaking and has the form (20). In our context the gauge kinetic function $\text{Ref}(A)$ is arbitrary and in principle contains non-renormalizable terms. In most of the models, again, one finds that the $|A|$ is of order M_P or larger, a fact that makes the non-renormalizable terms difficult to control similarly to the higher order terms in the K and W potential. Here, we will approximate $f(A)$ by polynomials and monomials or ascribe to it forms suggested by microscopic theories as the string theory.

In a FLRW background, neglecting spatial gradients, the Friedmann equation and the equation of motion for the ϕ (or the β) field are

$$H^2 = \frac{1}{3M_P^2} \left[\frac{\dot{\phi}^2}{2} (1 + 9M_*^{-2} H^2) + V(\phi) \right], \quad \partial_t \left[a^3 \dot{\phi} (1 + 3M_*^{-2} H^2) \right] = -a^3 V_\phi. \quad (24)$$

¹This formula is common for the old-minimal supergravity [32]. Nevertheless, a general supergravity-matter system in the new-minimal framework can be recast in this form after appropriate redefinitions [27].

Let us first demonstrate the advantages of the kinetic coupling to the inflationary applications. We assume here that the ϕ is the single inflating field and we consider a full polynomial potential without symmetry suppressed terms, actually non-renormalizable terms are naturally present in supergravity theories: $V(\phi) = \sum_n \lambda_n M_P^{4-n} \phi^n$. In the large field models of inflation the inflaton field has a value of order the Planck mass, M_P . This general potential cannot serve as large field inflationary model for the non-renormalizable terms, if not suppressed, spoil the flatness of the potential.

The non-minimal coupling of the kinetic term of the scalar field with the Einstein tensor G_{mn}

$$\mathcal{L} = -\frac{1}{2}\sqrt{-g} (g^{mn} - M_*^{-2}G^{mn}) \partial_m \phi \partial_n \phi \quad (25)$$

during a *de Sitter* phase takes the simple form $M_*^{-2}G^{mn} = -3M_*^{-2}H^2 g^{mn}$. For $HM_*^{-1} \gg 1$ the kinetic coupling implies that the canonically normalized scalar field is the $\tilde{\phi} = \sqrt{3}HM_*^{-1}\phi$. This rescaling recasts the polynomial potential in terms of the canonically normalized inflaton $\tilde{\phi}$ to the form

$$V(\tilde{\phi}) = \sum_n \lambda_n M_P^{4-n} \left(\frac{\tilde{\phi}}{\sqrt{3}HM_*^{-1}} \right)^n. \quad (26)$$

The non-renormalizable terms $\sum_{n=4}^{\infty} \lambda_n M_P^4 \left(\tilde{\phi} \times (\sqrt{3}HM_*^{-1}M_P)^{-1} \right)^n$ are suppressed by the “enhanced” mass scale $\sqrt{3}HM_*^{-1}M_P$. The slow roll parameters require $\tilde{\phi} > M_P$ and, hence, these higher order terms can be neglected and sufficient inflation can take place given that

$$M_P < \tilde{\phi} \ll M_P(HM_*^{-1}). \quad (27)$$

In terms of the field ϕ , which has non-canonical kinetic term, the above field-space region translates into

$$\frac{M_P}{HM_*^{-1}} < \phi \ll M_P. \quad (28)$$

This finding is of central importance since we work in a supersymmetric context. Even though we suggest a D-term inflation without a superpotential the generation of the inflationary potential is in principle not protected by any symmetry and in the most general case we cannot forbid the higher order terms.

We consider our theory as an effective one valid below some ultra-violet cut-off that we generally identify with the M_P . The field-space region (28) allows inflation to be realized in general form of potentials and reliable conclusions in this context can be derived. It can be said that the kinetic coupling theory is tailor-made for realizing an inflationary phase.

From a different perspective, if there is an internal symmetry that forbids the non-renormalizable terms and thereby suppresses the coefficients λ_n for $n \geq 5$ then inflation can be implemented in a much larger field-space region, than in the conventional (GR limit) large field inflationary models, that reads: $\phi > M_P/(HM_*^{-1})$

3.3 Introducing D-term inflationary potentials

We will attempt to capture some of the characteristics of the kinetic coupling in inflationary applications by considering some representative examples of inflationary potentials. We will concentrate on *single field inflation* models where one of the two fields is heavy enough and stabilized in the vacuum.

According to the Eqs. (24) and for $HM_*^{-1} \gg 1$ the slow-roll parameters of General Relativity (GR) $\epsilon \equiv M_P^2(V'/V)^2/2$ and $\eta \equiv M_P^2 V''/V$ are recast into

$$\tilde{\epsilon} \approx \frac{\epsilon}{3H^2 M_*^{-2}}, \quad \tilde{\eta} \approx \frac{\eta}{3H^2 M_*^{-2}}. \quad (29)$$

The requirement $\tilde{\epsilon}, |\tilde{\eta}| < 1$ yields that the field space region where slow-roll inflation is realized is rather increased. We will illustrate this below by considering different forms for the gauge kinetic function and thereby various types of potentials.

Linear potentials. Let us first assume that

$$f(A) = \frac{\xi^2}{2V_0} \sum_n \lambda_n \left(\frac{A}{M_P} \right)^n \quad (30)$$

where λ_n are real coefficients and we constrain the field to sub-Planckian values $|A| \ll M_P$. The scalar potential reads $\mathcal{V}(\phi, \beta) = V_0 (1 - \lambda_1 \phi/M_P - (\lambda_2 - \lambda_1^2) \phi^2/M_P^2 + \lambda_2 \beta^2/M_P^2 + \dots)$ where the ellipsis corresponds to negligible terms. The above potential includes two scalars that have a non-minimal derivative coupling. For $\lambda_1^2 \sim \lambda_2 > 0$ the ϕ field can be light enough and the β field can be heavy enough ($\gtrsim H$) and stabilized; hence the appearance of any sub-Planckian strong coupling scale [12] can be avoided. For $\phi \ll M_P$ the linear to ϕ term dominates and the potential reads:

$$\mathcal{V} \simeq V_0 \left(1 - \lambda \frac{\phi}{M_P} \right). \quad (31)$$

The slow-roll conditions yield the requirements for inflation $\tilde{\eta} = 0$ and $\tilde{\epsilon} \approx M_P^2 \lambda^2 / (2V_0 M_*^{-2}) < 1$.

Exponential potentials. If we now assume that the gauge kinetic function is of exponential form then we directly get an exponential type potential:

$$f(A) = \frac{\xi^2}{2V_0} e^{\lambda A/M_P} \quad \Rightarrow \quad \mathcal{V} = V_0 \frac{1}{\cos(\lambda \beta/M_P)} e^{-\lambda \phi/M_P}. \quad (32)$$

The form of the function $1/\cos x$ suggests that the β -dependent part of the potential will be stabilized with large enough mass to values $\langle 1/\cos(\lambda \beta/M_P) \rangle = 1$ and the ϕ will be the inflating field. An important reason for picking the gauge kinetic function (32) is that it is reminiscent of the dilaton coupling of string theory. Slow-roll inflation takes place for $\tilde{\eta} = 2\tilde{\epsilon} = \lambda^2/(3H^2 M_*^{-2}) < 1$ which corresponds to inflaton field values $\phi < M_P/\lambda \ln(V_0/M_*^2 M_P^2 \lambda^2)$, see Ref. [35] for further analysis.

Inverse power law potentials. As a final example, an inverse power law potential can be obtained if we consider monomial gauge kinematic function:

$$f(A) = \frac{\xi^2}{2V_0} \lambda \frac{A^n}{M_P^n} \quad \Rightarrow \quad \mathcal{V} = V_0 \frac{1}{\lambda} \frac{M_P^n}{\phi^n}, \quad (33)$$

where we used the relation $\text{Re}\{A^n\} = \phi^n \cos(n\theta)/(\cos\theta)^n$. We see that the field θ , the phase of the complex field $A = \rho e^{i\theta}$, is stabilized with large enough mass at $\theta = \kappa\pi$ and so $\text{Re}\{A^n\} = \phi^n$. Slow-roll inflation takes place for $\tilde{\epsilon} = (M_P^2 n^2)/(2\phi^2 \times 3H^2 M_*^{-2}) < 1$, $\tilde{\eta} = 2(1+n^{-1})\tilde{\epsilon} < 1$ which corresponds to inflaton field values $\phi^{n-2} < 2M_P^{n-4} V_0/M_*^2 \lambda n^2$.

For the above types of potentials an inflationary phase can be realized for a wider range of parameters. For the linear and the exponential, in GR limit, inflation is impossible for $\lambda \geq \mathcal{O}(1)$. It has to be $\lambda < \mathcal{O}(1)$ which implies, after absorbing λ to the mass scale, that the field ϕ has to be suppressed by a super-Planckian value. Here thanks to the kinetic coupling inflation is possible even for $\lambda \gtrsim 1$ and for sub-Planckian excursions for the (non-canonical) inflaton field ϕ .

3.4 Enhanced friction supergravity inflationary models and the CMB data

In order to make contact between the theory and observation the spectra of scalar and tensor perturbations have to be estimated [1, 36]. The density perturbations $\delta\rho$ of the inflaton are encoded in the

variable $\zeta = \delta\rho/(\rho+p)$ which is conserved on large scales in the absence of entropy perturbations and can be directly related to the cosmic microwave background temperature fluctuations. The power spectrum of the ζ variable in first order in the slow roll parameter $\tilde{\epsilon}$ reads [12] $\mathcal{P}_\zeta \approx H^2/8\pi^2\tilde{\epsilon}c_sM_P^2$. The sound speed squared having a dependence $c_s^2 \propto H^2\tilde{\epsilon}$, is subluminal and modifies the spectral tilt dependence on the slow roll parameters. The result is

$$n_s - 1 \equiv \left. \frac{d \ln \mathcal{P}_\zeta}{d \ln k} \right|_{c_s k = aH} \approx -8\tilde{\epsilon} + 2\tilde{\eta} \quad (34)$$

contrary to the well known GR limit formula $n_s - 1 = 2\eta - 6\epsilon$. The ratio of the tensor to scalar amplitudes, $r \equiv \mathcal{P}_g(k_*)/\mathcal{P}_\zeta(k_*)$, has the conventional GR dependence on the slow-roll parameter $\tilde{\epsilon}$ at the lowest order: $r = 16\tilde{\epsilon}$. However the new relation (34) allows for larger values for the $\tilde{\eta}$ slow-roll parameter. Namely it is

$$r = 2(1 - n_s) + 4\tilde{\eta} \quad (35)$$

instead of $r = (8/3)(1 - n_s) + (16/3)\eta$. Hence, given that $1 - n_s \cong 0.04$ and $r < 0.11$ [33], positive values for the $\tilde{\eta}$ can be accommodated, which correspond to potentials with $V'' \geq 0$. The number of e-folds, $N \equiv \int H dt$, for $H^2 M_*^{-2} \gg 1$ is given by the expression

$$N(\phi) = \frac{1}{M_P^2} \int_{\phi_f}^{\phi} (1 + 3H^2 M_*^{-2}) \frac{V}{V'} d\phi \approx \frac{1}{M_P^4} \int_{\phi_f}^{\phi} M_*^{-2} \frac{V^2}{V'} d\phi \quad (36)$$

The Planck Collaboration estimated the spectral index n_s from the observational data (Planck and WMAP) to be [33] $n_s = 0.9603 \pm 0.0073$ and the upper bound on the tensor to scalar ratio at $r < 0.11$. This constraint on r corresponds to an upper bound on the energy scale of inflation $H(\phi_*)/M_P \leq 3.7 \times 10^{-5}$ which implies that $\tilde{\epsilon}(\phi_*) \equiv \tilde{\epsilon}_* < 0.008$. On the other hand, the BICEP2 Collaboration [34] reported a value $r \sim 0.2$ which allows larger values for the $\tilde{\epsilon}_*$. The ϕ_* denotes the field value during inflation that the pivot scale $k_* = 0.002 \text{Mpc}^{-1}$ [33] exited the Hubble radius (not to be confused with the subscript at the mass scale M_* of the non-minimal derivative coupling).

Let us now examine the supergravity D-term inflationary models of the previous subsection in the light of the observational data. The key ingredient is that the inflaton field is characterized by the non-minimal derivative coupling to the Einstein tensor.

The *linear* model (31) yields a spectral index $1 - n_s = 8\tilde{\epsilon} = 4\lambda^2 M_P^2 M_*^2 / V_0$ which is related to the number of e-folds by the expression $1 - n_s \approx 4\lambda \Delta\phi / N(\phi_*) M_P$. For $N(\phi_*) \sim 50$ and $\Delta\phi \ll M_P$ this model can give a spectral index value $1 - n_s \sim 0.04$ and $r \sim 0.08$ for $\lambda \gg 1$. For the *exponential potential* (32) one finds $\tilde{\eta} = 2\tilde{\epsilon}$, $1 - n_s = 8\tilde{\epsilon} - 2\tilde{\eta} = 4\tilde{\epsilon}$ and $1 - n_s \sim 2/N(\phi_*)$. Hence, it predicts a tensor-to-scalar ratio $r = 0.16$ for $1 - n_s = 0.04$. This value of r lies between the Planck and the BICEP2 data. In the case of the *inverse power-law* models (33) the spectral index is given by the modified expression $1 - n_s = 4\tilde{\epsilon}(1 - n^{-1})$ and it is in tension with the Planck data however, it is viable according to the BICEP2 data for large enough power n .

The kinetic coupling operates like an enhanced friction and inflation takes place more generically than in the GR limit. Since inflation is primarily introduced to address (or better ameliorate) the homogeneity, isotropy and flatness problem we can say that even excluded models are still motivated candidates for inflation in the context of supergravity kinetic coupling. Afterwards, in order to seed the large-scale structure formation in the universe, a mechanism like the curvaton or the modulated reheating may take place in the post-inflationary universe.

The non-minimal derivative coupling of the scalar field ϕ with the Einstein tensor in supergravity renders the ϕ a compelling inflaton candidate due to both the enhanced friction effect and the symmetry

suppression of the F-terms. Despite these advantages the absence of tree level superpotential interactions may be problematic for a sufficient reheating of the universe. However, we note that there is the coupling of the inflating scalar ϕ with the gauge field strength (19) and, also, interactions beyond the tree level. Moreover, there are mechanisms like the gravitational particle production [38, 39] invented for such questionable situations.

4 Conclusions

In the present paper we have examined the implementation of an inflationary phase by the scalar component of a chiral superfield in a supergravity context. The particular characteristic of this scalar is that it is non-minimally coupled to the Einstein tensor G^{mn} , and the slow-roll conditions can be satisfied more generically. The potential is introduced via a Fayet-Iliopoulos D-term since the R-symmetry of the theory excludes the introduction of a superpotential, and its gauging is also forbidden due to stability issues. Even though this is a higher derivative theory it does not give rise to ghost instabilities.

The model we propose is a pure D-term inflation with a gravitational enhanced slow-roll for the inflaton. These two features have important implications for the inflationary dynamics. Firstly, the accelerated expansion can be realized for sub-Planckian excursions for the (non-canonical) inflaton field as well as for sub-Planckian parameter scales. Sub-Planckian excursions for the inflaton field are welcome because possible non-renormalizable terms are suppressed. This fact together with the absence of the F-terms render this model free from the notorious η -problem of supergravity. Secondly, the field space region where the slow-roll conditions are satisfied is increased. Hence an inflationary period is realized for more generic initial conditions. Thirdly, the relation between the spectral index of the scalar perturbations and the slow-roll parameters is modified due to the corrected sound speed of the scalar perturbations. These imply that some inflationary models may provide a better fit to the data or even render some excluded inflationary models observationally viable. For example, here, the predictions of the exponential potential can be accommodated by the combined Planck and BICEP2 data [35].

Concluding, we comment that throughout this work we have considered single field inflationary potentials. Although their inflationary dynamics are influenced by the value of the new scale M_* , these potentials are characterized only by Planck mass suppression scale.

We believe that this work has offered some different insight into how inflation might work in a supergravity framework.

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